

# Kaluza-Klein Gluons as a Diagnostic of Warped Models

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## Abstract

We study the properties of  $g^1$ , the first excited state of the gluon in representative variants of the Randall Sundrum model with the Standard Model fields in the bulk. We find that measurements of the coupling to light quarks (from the inclusive cross-section for  $pp \rightarrow g^1 \rightarrow t\bar{t}$ ), the coupling to bottom quarks (from the rate of  $pp \rightarrow g^1 b$ ), as well as the overall width, can provide powerful discriminants between the models. In models with large brane kinetic terms, the  $g^1$  resonance can even potentially be discovered decaying into dijets against the large QCD background. We also derive bounds based on existing Tevatron searches for resonant  $t\bar{t}$  production and find that they require  $M_{g^1} \gtrsim 950$  GeV. In addition we explore the pattern of interference between the  $g^1$  signal and the non-resonant SM background, defining an asymmetry parameter for the invariant mass distribution. The interference probes the relative signs of the couplings of the  $g^1$  to light quark pairs and to  $t\bar{t}$ , and thus provides an indication that the top is localized on the other side of the extra dimension from the light quarks, as is typical in the RS framework.

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## I. INTRODUCTION

The large hierarchy between the Planck scale where quantum gravity effects are important, and the scale where the electroweak symmetry is broken, drives the wealth of models at the electroweak scale, and motivates the Large Hadron Collider (LHC) experiments. While weakly coupled supersymmetry remains a leading candidate to stabilize the hierarchy, the Randall-Sundrum (RS) models of a warped extra dimension [1] have recently emerged as a fascinating alternative, which may be connected to string landscape solutions of the cosmological constant problem [2], and possess an interesting four dimensional dual interpretation in terms of the composite states of a strongly coupled nearly conformal field theory (CFT) [3].

The original RS model had all of the Standard Model confined to the IR brane (appearing as composites in the dual description). However, the RS solution to the hierarchy problem requires only the Higgs to be localized at the IR boundary, and there are compelling reasons to consider most of the SM might actually lie near the UV brane (and thus mostly fundamental with respect to the CFT in the dual description). Theories with the SM in the bulk can incorporate Grand Unification of couplings [4], motivate the flavor hierarchy of fermion masses [5], and incorporate a dark matter candidate [6]. However, such theories face significant challenge from precision electroweak observables [7], requiring specific features [8, 9, 10] in order to remain natural.

At the LHC, production of colored states is usually dominant, and the Kaluza-Klein (KK) excitations of the gluons are particularly attractive, because they are singly produced and thus have larger rates than the KK quarks. Thus, they are usually considered to be likely to be the first signs of warped physics, and the first excitation of the gluon ( $g^1$ ) the state for which we will have the most statistics available in order to unravel the details of the underlying theory. They are the natural place to explore whether or not we can use LHC data to determine which particular detailed RS model has been realized in nature, and which parameters describe it. Recently, significant work has begun on some of the simplest RS constructions to study the production and decay of the first KK mode of the gluon, in order to determine the reach of the LHC to discover RS through

its detection [11].

While the KK gluon is the most promising avenue to discover RS, it is nevertheless challenging. The coupling to the light quarks that are the primary constituent of the proton, while characterized by the strong coupling, are somewhat suppressed by the localization of the light fermions close to the UV boundary (in the CFT language, the light fermions are largely fundamental fields and couple to the gluon largely through a small mixing with CFT states). This leads to somewhat smaller production cross sections than are typical of QCD. The decay of the gluon is expected to be predominantly into top quarks, a consequence of the large top mass, which necessitates that top is itself located close to the IR brane (mostly composite). The tops are produced from a very heavy resonance, and are highly boosted, which makes it experimentally challenging to reconstruct them from the large QCD backgrounds [11].

In this article, we explore several more of the commonly considered theories which attempt to render RS consistent with precision electroweak data. We consider the model with a simple  $SU(2)$  [8] custodial symmetry (already studied before [11]) as a beginning, and also consider models with large brane kinetic terms [9] or an expanded custodial symmetry which protects the  $Z$ - $b\bar{b}$  vertex from large corrections [10, 12, 13] in order to characterize the difference in the properties of the first KK mode of the gluon in each case. We find that there are general features which can discriminate between the cases, and thus that the specific realization of the RS model leaves an imprint in the properties of the KK gluon.

The article is laid out as follows: in section II, we review the specific details of the models under consideration. In Section III we show the  $g^1$  production rates and decay properties and show how the strong coupling can lead to interesting finite width effects in section IV. Section V contains our conclusions.

## II. MODELS

### A. The Basic RS Model with the SM in the Bulk

The basic RS model is a slice of  $\text{AdS}_5$  with the background metric

$$ds^2 = \left(\frac{z_h}{z}\right)^2 [\eta_{\mu\nu} dx^\mu dx^\nu + (dz)^2], \quad (1)$$

with curvature  $\kappa = 1/z_h \lesssim M_{Pl}$ .  $x^\mu$  are the coordinates of the four large dimensions,  $z$  parameterizes the coordinate along the extra dimension, and  $\eta_{\mu\nu} = \text{Diag}(-, +, +, +)$  is the four-dimensional metric. Greek letters denote the four large dimensions 0, 1, 2, 3 and capital roman letters include the fifth dimension as well. The UV boundary is at  $z_h = 1/\kappa$  where the scale factor  $(z_h/z)^2 = 1$  and the IR boundary is at  $z_v \sim 1/\text{TeV}$ , as motivated by the hierarchy problem.

We are particularly interested in a model where all Standard Model (SM) fields, except perhaps the Higgs, propagate in the entire 5-d spacetime, and will be primarily concerned with the gluon and colored fermion fields. The action for the gauge fields and fermions is,

$$S = \int d^5x \sqrt{-g} \left\{ -\frac{1}{4g_5^2} F_{MN}^a F^{MN a} + i\bar{\Psi} \Gamma^{\dot{M}} e_M^{\dot{M}} D_M \Psi + i c \kappa \bar{\Psi} \Psi \right\}. \quad (2)$$

where  $\Gamma^{\dot{M}}$  are the 5d ( $4 \times 4$ ) Dirac matrices,  $e_M^{\dot{M}}$  is the vielbein,  $a$  is an adjoint gauge index and  $c$  parameterizes the magnitude of a bulk mass for the fermion in units of the curvature.

We work in a unitary gauge  $A_5 = 0$ , and decompose the 5d fields in KK modes,

$$A_\mu^a(x, z) = \sum_n A_\mu^{a(n)}(x) g^{(n)}(z), \quad (3)$$

$$\Psi_{L,R}(x, z) = (\kappa z)^{3/2} \sum_n \psi_{L,R}^n(x) \xi_{L,R}^{(n)}(z). \quad (4)$$

The wave functions are given by combinations of Bessel functions

$$g^{(n)}(z) = N_n z [J_1(m_n z) + b_n Y_1(m_n z)]. \quad (5)$$

with normalization factor  $N_n$  and admixture controlled by  $b_n$ . The mass spectrum is controlled by the boundary conditions, with the masses satisfying,

$$b_n = -\frac{J_0(m_n z_h)}{Y_0(m_n z_h)} = -\frac{J_0(m_n z_v)}{Y_0(m_n z_v)}. \quad (6)$$

For an unbroken gauge group, there is a zero mode with wave function  $g^0(z) = 1/\sqrt{L}$ ,  $L = 1/k \log z_v/z_h$ . Of particular note for the following is the fact that the light KK states have most of their support close to the IR boundary.

The physics of bulk fermions was worked out in [14]. The spectrum depends sensitively on the bulk mass term  $c$ . To remove unwanted light degrees of freedom, we impose the boundary conditions such that either the right- or the left-chiral zero mass component is projected out. The KK states form left- and right-chiral pairs whose wave functions are also Bessel functions,

$$\xi_{\pm}^{(n)}(z) = \mathcal{N}_n(\kappa z) [J_{|c \pm 1/2|}(m_n z) + \beta_n Y_{|c \pm 1/2|}(m_n z)] , \quad (7)$$

where  $- (+)$  are for the right- (left-) chiral modes, and the masses are determined by imposing the equality,

$$\beta_n = \frac{J_{|c-1/2|}(m_n z_h)}{Y_{|c-1/2|}(m_n z_h)} = \frac{J_{|c-1/2|}(m_n z_v)}{Y_{|c-1/2|}(m_n z_v)} , \quad (8)$$

and  $\mathcal{N}_n$  is a normalization factor. The zero mode wave functions are,

$$\xi^{(0)}(z) = \mathcal{N}_0 (\kappa z)^{1/2-c_\Psi} . \quad (9)$$

These wavefunctions assume the right-handed zero mode is the one allowed by the boundary conditions; the left-handed case is given by  $c \rightarrow -c$ . The zero mode is exponentially peaked toward the UV boundary for  $c < -1/2$  and toward the IR for  $c > -1/2$ . To avoid confusion, we adopt a notation where  $c$ 's explicitly refer to right-chiral fields, so the left-chiral fermions should be understood to actually have  $-c$  as their mass parameter.

Assuming  $\mathcal{O}(1)$  5D Yukawa couplings, the hierarchy in the effective 4D Yukawa couplings can be motivated by the exponential suppression of the wave functions at the IR boundary for order one differences in  $c$ . In particular, one cannot allow strong suppression of the top quark wave functions on the IR boundary, because to reproduce the observed top mass one would have to adjust the 5d Yukawa coupling to be too strong to have a perturbative description. There is further motivation from precision electroweak data [8, 9], which prefers the light fermion mass parameters to be close to  $-1/2$  (including the left-handed top, as it comes along with the left-handed bottom, leading to tension in the

choice of  $c$  for  $Q_3$ ) in order to cancel the leading contribution to the  $S$  parameter from the weak boson KK modes. With this setup, and the additional suppression of the  $T$  parameter from a custodial  $SU(2)$ , masses of the KK modes of around 3 TeV are roughly consistent with precision measurements.

Specifically, we consider  $c_{t_R} \sim 0$ ,  $c_{Q_{3L}} \sim 0.4$ , and all others  $c_f \lesssim -0.5$ . As we will see shortly, the physics we study does not depend strongly on  $c$  once it is  $< -1/2$ , so the specific values for light fermions are not important. The choices of  $c$  specify the fermion zero mode wave functions, and we compute the couplings of the first KK gluon to the fermion zero modes as the integral over the wave functions. The light quarks all have very similar couplings of roughly  $g_f \simeq -g_S/5$ , the third family left-handed quarks  $g_{Q3} \simeq g_S$ , and the right-handed top quark  $g_t \simeq 4g_S$ , where  $g_S$  is the strong coupling constant which characterizes the coupling of the gluon zero mode.

## B. IR Brane Kinetic Terms

An alternate way to render precision electroweak data consistent with low KK mode masses is to include large-ish kinetic terms for the gauge fields on the IR brane [9]. Such terms repel the KK mode wave functions from the brane, and have a large effect on the phenomenology of the KK modes [15]. Brane terms are a class of higher dimensional operators of the 5d theory,

$$-\frac{1}{4g_5^2} \int d^5x \sqrt{-g} \{ F_{MN}^a F^{MN a} \} 2 r_{IR} \delta(z - z_v) \quad (10)$$

and will be induced by the orbifold boundary conditions and localized fields [16]. Their magnitude  $r_{IR}$  is a free parameter of the effective theory. While the size of the IR boundary kinetic term for the gluon is not closely connected to the quality of the electroweak fit, one would expect that if the UV physics is such that there are large IR kinetic terms for the electroweak bosons, such terms are probably also large for the gluon as well. Thus, if one could infer the presence of large gluon terms, it would at least suggest that a similar term is present in the electroweak sector, and responsible for the success of the SM in explaining the electroweak fit.

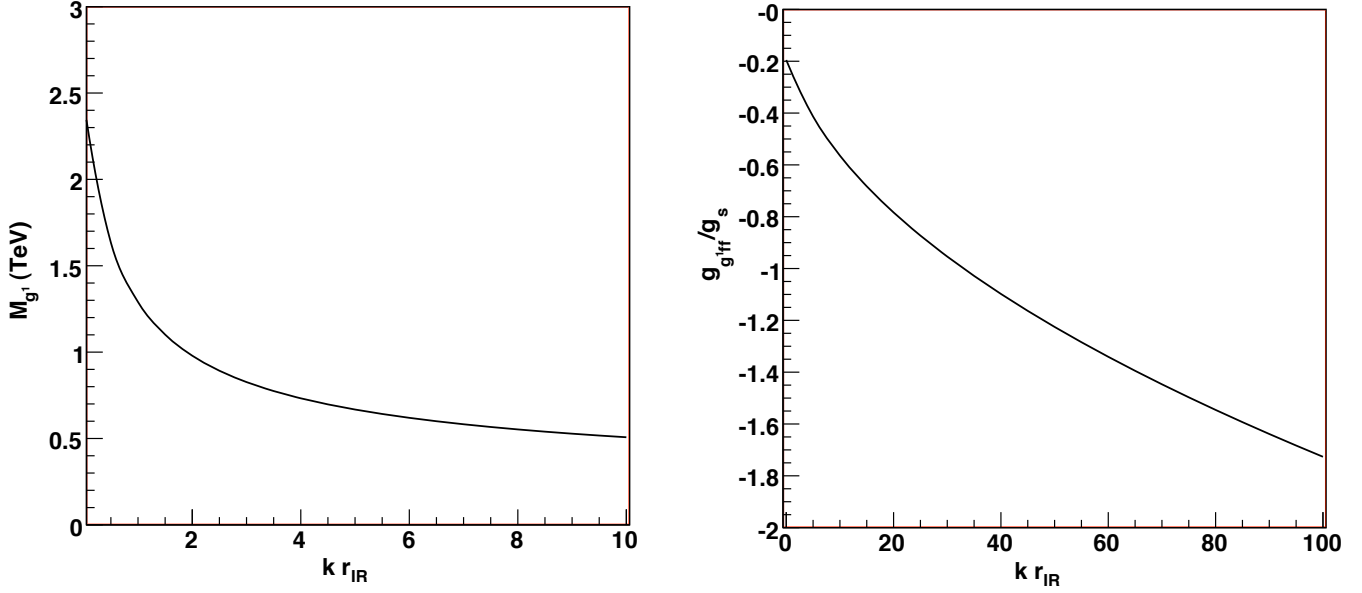


FIG. 1: The 1st KK gluon mass in units of  $1/z_v$  and coupling of the first KK gluon to a fermion zero mode localized at UV brane as a function of brane kinetic term  $\kappa r_{IR}$ .

The IR boundary kinetic term does not affect the form of the bulk wave functions, Eq. (5). The boundary conditions become

$$b = -\frac{J_0(m_n z_h)}{Y_0(m_n z_h)} = -\frac{J_0(m_n z_v) - (\kappa r_{IR}) m_n z_v J_1(m_n z_v)}{Y_0(m_n z_v) - (\kappa r_{IR}) m_n z_v Y_1(m_n z_v)}, \quad (11)$$

indicating different admixture of the Bessel functions  $J_1$  and  $Y_1$  in the solutions. While there is no analytic solution for the masses, they may be easily obtained numerically. In Figure 1, we show the variation of the first KK mode gluon mass and coupling to UV-localized states as a function of the magnitude of the IR brane term  $\kappa r_{IR}$ . In Figure 2, we show the dependence of the coupling on  $c$  for a few different choices of  $\kappa r_{IR}$ . The inclusion of the boundary terms ameliorates the strongest constraints from precision electroweak data, and opens up considerably more freedom to choose the fermion  $c$ 's. However, in computing properties below, we imagine a situation in which the  $c$ 's are as in the  $SU(2)$  custodial version outlined above (for example, to explain the flavor hierarchies), with large contributions to the electroweak  $T$ -parameter controlled by the IR boundary kinetic terms.

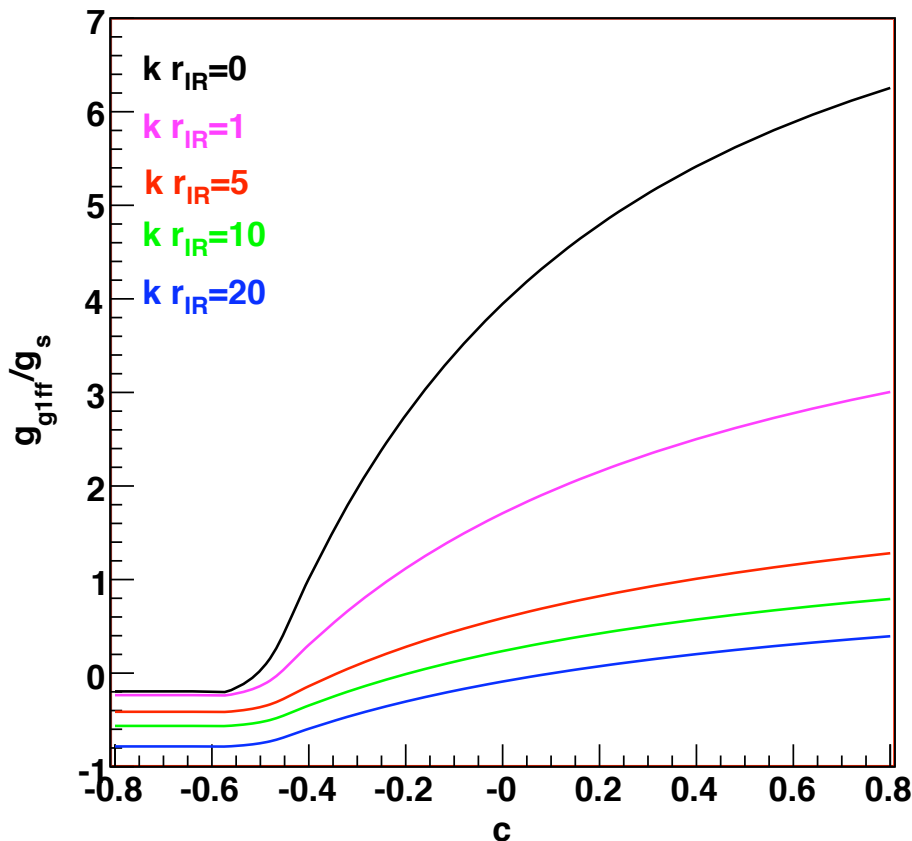


FIG. 2: Coupling of the first KK gluon (with respect to the zero mode gluon coupling) with  $\kappa r_{IR} = 0, 1, 5, 10, 20$  (descending) to a fermion zero mode as a function of bulk mass parameter  $c$ .

### C. Holographic Higgs with Expanded Custodial Symmetry

The models with a custodial  $SU(2)$  symmetry or large IR boundary kinetic terms (combined with the choices of the  $c$ 's motivated above) continue to be challenged by the large top mass, which we saw did not allow  $Q_3$  to be pushed quite as far away as was optimal for the lighter fermions. This results in corrections to the  $Z$ - $b_L$ - $\bar{b}_L$  coupling compared to those of light fermions which are slightly too large for the experimental errors, and push in a direction unhelpful for  $A_b^{FB}$  [17].

In [10], it was noticed that a subgroup of the custodial symmetry can protect the



$Z$ - $b_L$ - $\bar{b}_L$  coupling, provided the third generation doublet is embedded in a representation for which the  $SU(2)_L$  and  $SU(2)_R$  representations (and the third component of each) are the same. This implies that to better fit  $Z$ - $b_L$ - $\bar{b}_L$ , we expand  $Q_3$  into a bi-doublet under  $(SU(2)_L, SU(2)_R)$ . The unwanted additional fermions in the bi-doublet are removed from the zero mode spectrum by adjusting their boundary conditions. Having promoted  $Q_3$  to a bi-doublet, we recover freedom to consider the  $c$  parameter for  $Q_3$  very different from  $-1/2$ .

In order to have a specific framework, we analyze the model of gauge-Higgs unification [13] (similar to an earlier model [18]) in which the allowed parameter space is analyzed in great detail [19], reproducing light fermion masses and mixings, and demanding consistency with flavor-changing neutral currents induced by the KK modes of the gauge bosons. While some of the features are particular to the gauge-Higgs unified model and the mechanism by which it realizes fermion masses and mixings, some of the most important features are fairly generic to models in which an expanded custodial symmetry is protecting  $Z$ - $b_L$ - $\bar{b}_L$ .

The bulk gauge symmetry is  $SO(5) \times U(1)_X$ , broken by boundary conditions to  $SU(2)_L \times SU(2)_R \times U(1)_X$  on the IR boundary, and to the Standard Model  $SU(2)_L \times U(1)_Y$  gauge group on the UV brane [18]. The  $U(1)_X$  charges are adjusted so as to recover the correct hypercharges, where  $Y/2 = T_R^3 + Q_X$  with  $T_R^3$  the third  $SU(2)_R$  generator and  $Q_X$  the  $U(1)_X$  charge. As motivated above, we wish  $Q_3$  to be part of a bi-doublet, and an economical choice is to embed it in a  $\mathbf{5}_{2/3}$  of  $SO(5)$  (the subscript refers to the  $U(1)_X$  charge). As discussed in [19], it is preferable to place  $t_R$  in a separate  $\mathbf{5}_{2/3}$  to avoid large negative corrections to the  $T$  parameter.  $b_R$  is part of a  $\mathbf{10}_{2/3}$ , allowing for the bottom Yukawa coupling, and the first and second generations are replicas of this structure in order to generate CKM mixing in a straight-forward way. Enhanced coupling to bottom quarks is also potentially a signal of RS attempts to explain the observed deviation in  $A_{FB}^b$  [20].

The scan over parameters of [13] prefers that the quarks and leptons of the first two generations are localized close to the Planck boundary in order to suppress flavor changing neutral currents. The expanded custodial symmetry, combined with relatively light KK

modes for the  $Q_3$  custodial partners, is so efficient at suppressing contributions to the  $T$  parameter, that it reduces some of the usual SM top contribution, and can result in  $T$  large and negative, in conflict with the electroweak fit [21]. To ameliorate this new concern, the freedom to consider  $Q_3$  closer to the IR boundary (compensated by moving  $t_R$  somewhat away from it) is crucial, allowing  $c_{Q_{3L}} \sim 0.2$ ,  $c_{t_R} \sim -0.49$ , and  $c_f \lesssim -0.5$ , for which the couplings to the first KK mode of the gluon are approximately  $g_f \sim -g_S/5$ ,  $g_t \sim 0.07g_S$ , and  $g_{Q_3} \sim 2.76g_S$ .

This model generically leads to very light KK quarks, the lightest of which are the  $SO(5)$  bi-doublet partners of the right-handed up-type quarks of the first two generations  $u_i$  (by virtue of the choice of  $c$  for the two light generations) [13]. Each generation contains

$$Q_{2R}^i = \begin{pmatrix} \chi_{2R}^{u_i}(+, -) & q_R^{u_i}(+, -) \\ \chi_{2R}^{d_i}(+, -) & q_R^{d_i}(+, -) \end{pmatrix}, \quad (12)$$

along with their  $(-, +)$  left-handed Dirac partners. The  $(\pm, \mp)$  refers to their boundary conditions on the (UV, IR) boundaries, and do not lead to zero modes (as desired), and modify the equation which determines their masses and admixture of Bessel functions. For the right-handed  $(+, -)$  states this leads to,

$$\beta_n = \frac{J_{|c-1/2|}(m_n z_v)}{Y_{|c-1/2|}(m_n z_v)} = \frac{J_{|c-1/2|\mp 1}(m_n z_h)}{Y_{|c-1/2|\mp 1}(m_n z_h)}, \quad (13)$$

with upper(lower) signs for  $c > -1/2$  ( $c < -1/2$ ). The left-handed  $(-, +)$  states satisfy,

$$\beta_n = \frac{J_{|c+1/2|}(m_n z_h)}{Y_{|c+1/2|}(m_n z_h)} = \frac{J_{|c+1/2|\mp 1}(m_n z_v)}{Y_{|c+1/2|\mp 1}(m_n z_v)}. \quad (14)$$

Armed with these wave functions, we compute the coupling of these potentially light first KK modes of the custodial partners to the first KK mode of the gluon. The results for both chiralities are presented in Figure 3, and indicate that one chirality is always very strongly coupled,  $g \sim 6g_S$ , irrespective of the value of  $c$ .

## D. A Warped Higgsless Model

A final variant of the warped theory has no Higgs, and breaks the electroweak symmetry by boundary conditions [22]. The need for the KK modes of the weak vector bosons to

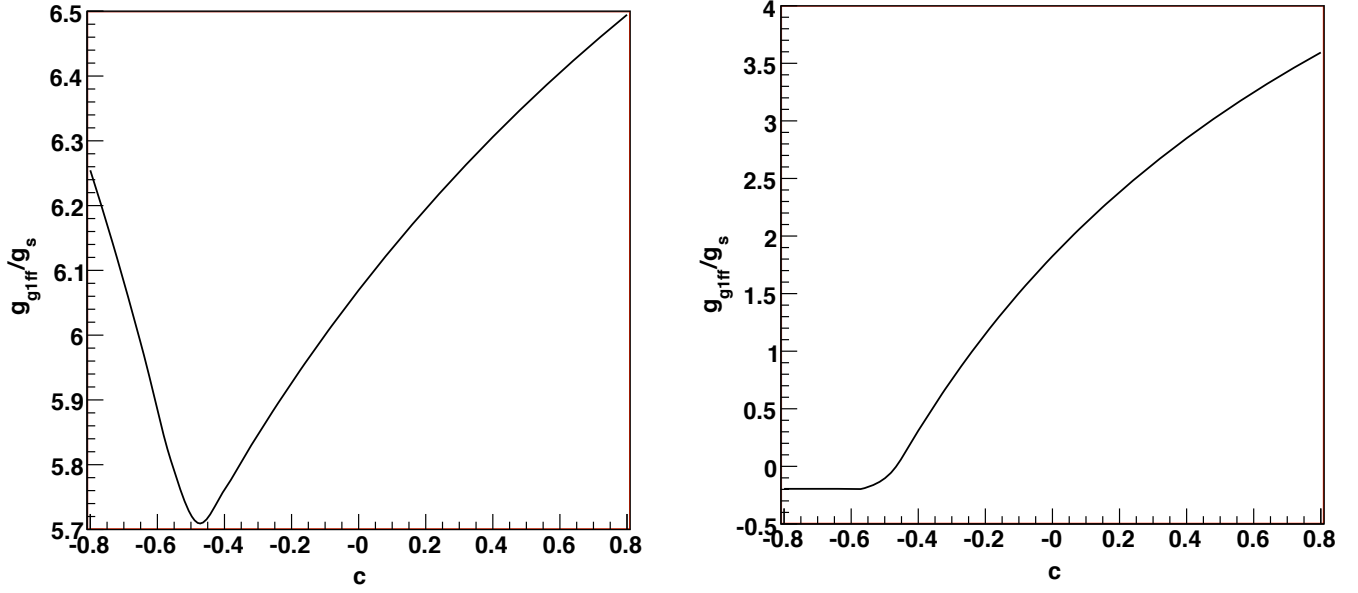


FIG. 3: Coupling of the First KK mode of the gluon to the light KK modes of the custodial partners of the right-handed up-type quark as a function of the bulk mass parameter  $c$ . The left panel shows the left-handed coupling whereas the right panel the right-handed coupling.

unitarize  $WW$  scattering implies that the scale of KK mode masses is at most several hundred GeV, whereas the need to be consistent with precision electroweak data and realize a large top mass requires [23][27]

$$\begin{aligned} g_t &= 2.5g_S, \quad g_{Q3} = 2g_S, \quad g_b = -0.32g_S, \\ g_{\text{otherRH}} &= -0.33g_S, \quad g_{\text{otherLH}} = 0.15g_S. \end{aligned} \quad (15)$$

We see that the basic trend is very similar to the other RS models we consider. The main distinguishing feature is the fact that the mass of the KK gluon must be less in order for the Higgsless model to remain consistent with perturbative unitarity.

### III. PRODUCTION AND DECAY

The details of production of KK gluons at the LHC will depend on how they couple to the relevant partons at LHC energies, and these differences will give us a powerful way

to discriminate between models. Note that the vertex with two gluons and a KK gluon is zero at tree-level, meaning that the dominant production mode is  $q\bar{q}$  annihilation. As is well-known, in the standard RS framework the KK gluon coupling to all fermions aside from  $t_R$  is suppressed. As we saw above, models with brane kinetic terms can increase couplings to UV-localized fields, which increases the rate and affect the branching ratios. In addition, the models with custodial symmetry have a large coupling to  $Q_3$ , turning on a new production mode from bottom fusion, but have a smaller branching ratio because the decay into the custodial partner KK modes may compete with top. In Figure 4, we plot the cross section, calculated at leading order by MADGRAPH [25],  $pp \rightarrow g^1$  for  $\sqrt{S} = 14$  TeV, as a function of  $g^1$  mass for the models considered above, including the channels initiated by light quark fusion and bottom fusion.

As indicated above, models with the extra custodial symmetry to protect the  $Z$ - $b_L$ - $b_L$  coupling from large corrections have considerably more freedom to locate  $Q_3$  closer to the IR brane, and considerations of the  $T$  parameter prefer to do so. This enhances the  $g^1$  coupling to left-handed bottoms (up to about  $3g_S$ ) and results in large production from bottom quark fusion, as shown in Figure 4. It would be useful to be able to discern that the increase over the expected production rate in the standard bulk SM RS picture is because of this enhancement of the coupling to bottom (which would be suggestive of the expanded custodial symmetry), as opposed to a straight enhancement of the coupling to all light quarks (which would be more suggestive of a large kinetic term on the IR boundary). One could study the rapidity distribution of the  $g^1$  itself (as reflected in the final state top pair distribution). The fact that both  $b$  and  $\bar{b}$  are sea quarks would imply a more central rapidity distribution than would result from  $q$  and  $\bar{q}$ , because  $q$  as a valence quark will tend to carry more momentum than  $\bar{q}$ . However, the  $g^1$  rapidity distribution is only modestly sensitive to the initial state, and is also sensitive to the  $g^1$  mass and width. Thus, we turn to a more straight-forward measure of the contribution of  $b\bar{b}$  to  $g^1$  production [28] which is to compare the rate of  $g^1 \rightarrow t\bar{t}$  to  $bg^1 \rightarrow bt\bar{t}$ . In Figure 5, we present these rates for standard RS, the model with  $\kappa r_{IR} = 5$ , and the model with  $Q_3$  localized around the IR brane. We find that as expected, the rate for the model with custodial symmetry is enhanced by the large bottom coupling by about an order of

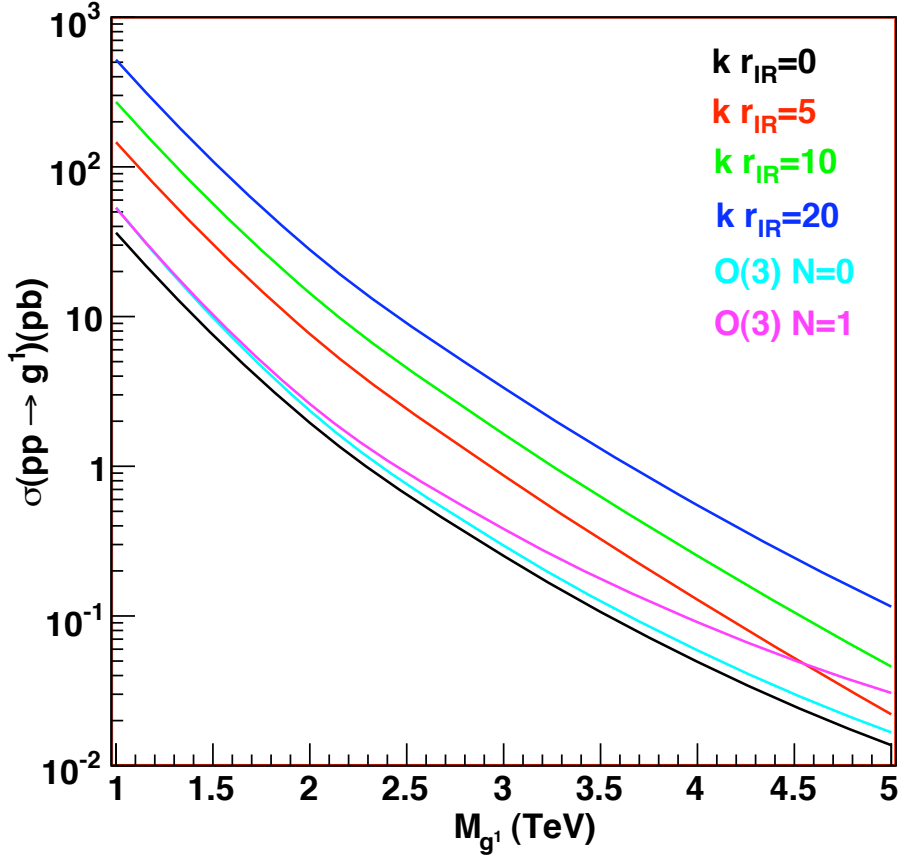


FIG. 4: Cross section for  $pp \rightarrow g^1$  at the LHC, for standard RS with the SM in the bulk ( $\kappa r_{IR} = 0$ ), three models with large brane kinetic terms ( $\kappa r_{IR} = 5, 10, 20$ ) and the model with a larger custodial symmetry, in the cases when  $N = 0$  or  $1$ , of the additional KK custodial partner quarks are light enough that  $g^1$  can decay into them.

magnitude. In addition, the model with IR boundary kinetic terms shows a rate which is suppressed by a factor of about five, because while the boundary kinetic term slightly enhances the coupling of the UV-localized  $b_R$ , it more dramatically suppresses the coupling to the IR-localized  $b_L$  (c.f. Figure 2). Ultimately, one must include the SM background and detector efficiencies for a specific decay channel of  $g^1$ . As a step in this direction, in Figure 6 we plot the differential cross-section for both the  $pp \rightarrow t\bar{t}$  and  $pp \rightarrow b\bar{t}t$  signals and SM backgrounds with respect to the  $t\bar{t}$  invariant mass, in the standard RS model

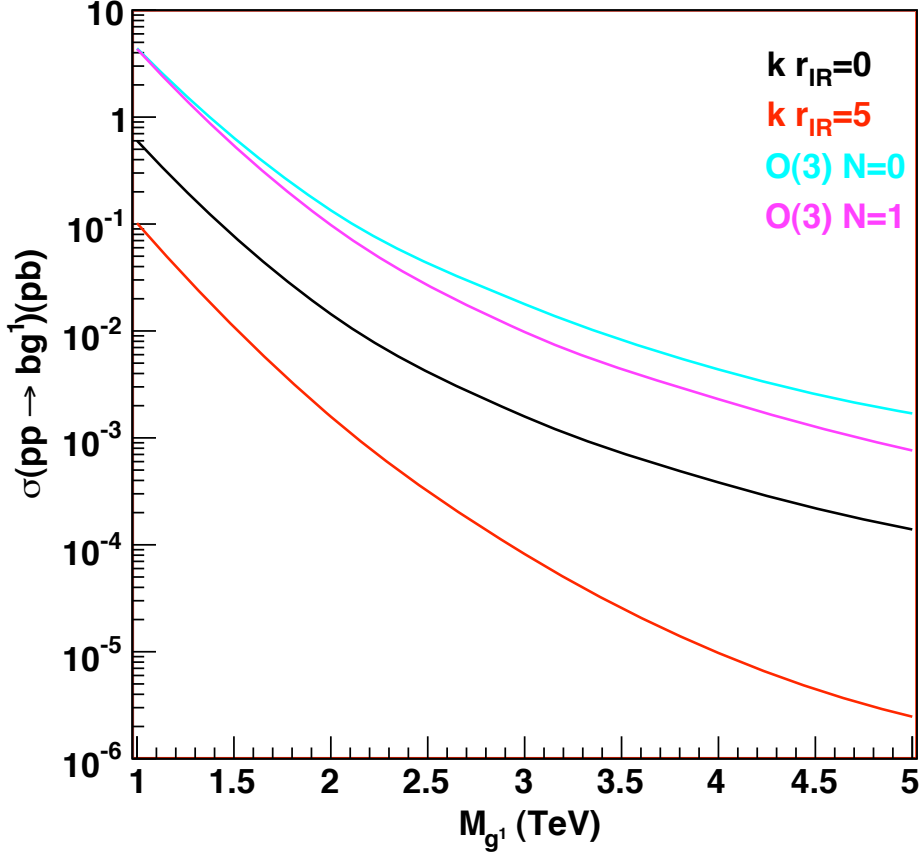


FIG. 5: Cross section for  $pp \rightarrow bg^1$  at the LHC, for standard RS with the SM in the bulk ( $\kappa r_{IR} = 0$ ), a model with a large brane kinetic term ( $\kappa r_{IR} = 5$ ) and the model with a larger custodial symmetry, in the cases when 0 (1) of the additional KK custodial partner quarks are light enough that  $g^1$  can decay into them.

and one with a larger custodial symmetry. In both cases, for  $M_{g^1} = 2$  TeV, a peak is visible above the SM background, and the size of  $g^1 b$  production relative to  $g^1$  production discriminates between the two models.

The width of  $g^1$  is strongly dominated by the states close to the IR brane to which it couples strongly. Generically, the partial width into  $f\bar{f}$  for which the left- and right-chiral

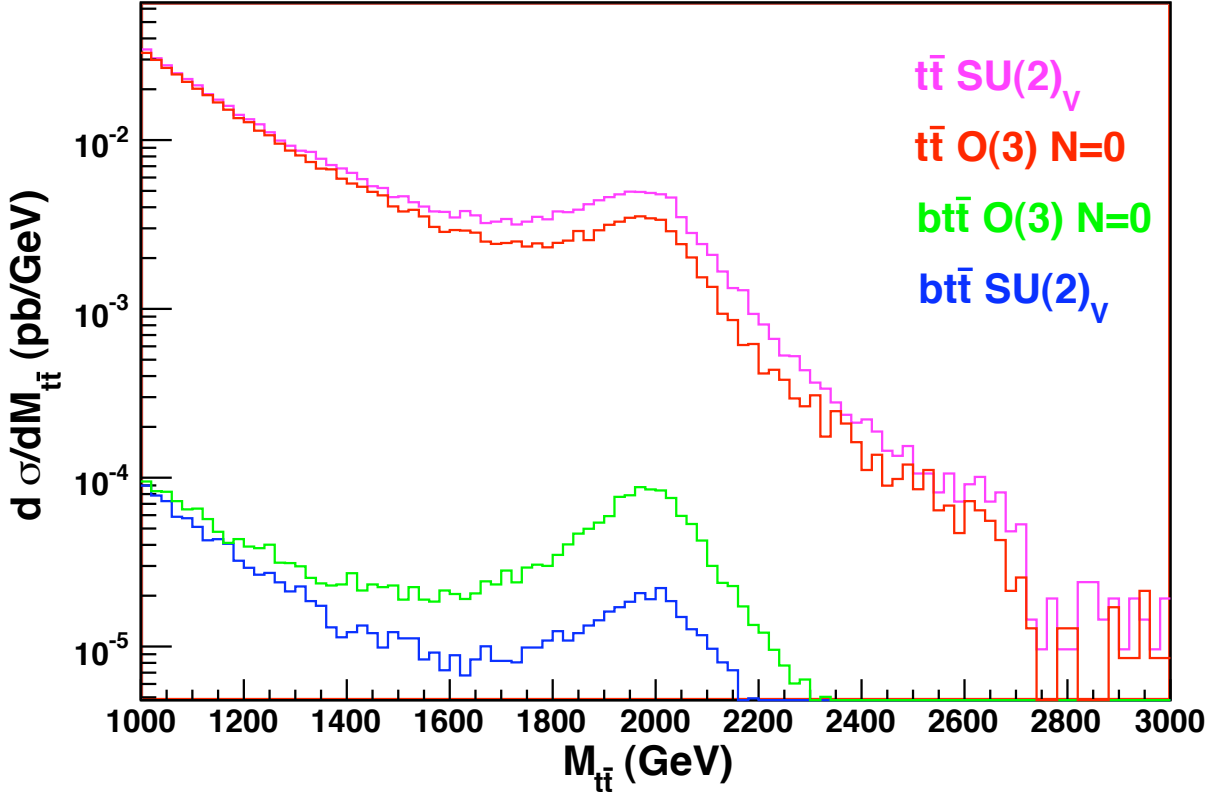


FIG. 6: Invariant mass distribution of  $t\bar{t}$  in the standard RS model ( $SU(2)_V$  custodial symmetry) and the model with a larger  $O(3)$  custodial symmetry in  $pp \rightarrow t\bar{t}$  and  $pp \rightarrow b\bar{t}\bar{t}$  respectively.

interactions with  $g^1$  are  $g_L$  and  $g_R$  is given by,

$$\begin{aligned} \Gamma_{G1 \rightarrow f\bar{f}} &= \frac{1}{48\pi M_{g^1}} \sqrt{1 - \frac{4m_f^2}{M_{g^1}^2}} [(g_L^2 + g_R^2) (M_{g^1}^2 - m_f^2) + 6g_L g_R m_f^2] \\ &\simeq \frac{M_{g^1}}{48\pi} (g_L^2 + g_R^2) , \end{aligned} \quad (16)$$

where the final approximation holds in the limit  $M_{g^1} \gg m_f$ . Decays to top quarks are always important, because either  $t_R$  or  $Q_3$  must be IR-localized to realize the large top mass. In addition, when the custodial partner KK quarks are light enough for  $g^1$  to decay into them, they will also take a substantial fraction of the branching ratio, because they are also IR-localized and have large coupling. The IR boundary kinetic terms can suppress the coupling to top, and enhance the decay into light quarks. In Table I we list

Model	top quarks	bottom quarks	light quarks	custodial partners	$\Gamma_{g^1}/M_{g^1}$
Basic RS	92.6%	5.7%	1.7%		0.14
$\kappa r_{IR} = 5$	2.6%	13.2%	84.2%		0.11
$\kappa r_{IR} = 20$	7.8%	15.1%	77.1%		0.05
$O(3), N = 0$	48.8%	49.0%	2.0%		0.11
$O(3), N = 1$	14.6%	14.6%	0.6%	70.2%	0.40

TABLE I: The branching ratios of  $g^1$  into tops, bottoms, light quarks (jets), and custodial partners, as well as the total width  $\Gamma_{g^1}/M_{g^1}$ , for several different RS scenarios in the limit  $M_{g^1} \gg m_f$ .

the branching ratios into top quarks, bottom quarks, light quarks (jets) and exotic quarks in several different RS models. The total width also sensitively depends on the couplings, and how many custodial partners are available as decay modes. The width is generally large, owing to the strong couplings present, and it may be possible to reconstruct it from the final state invariant mass distributions, which would also allow one to use it as an additional source of information. The final column of Table I shows the total width  $\Gamma_{g^1}/M_{g^1}$  for each model. Variations are typically around 5%, with the exception of the model with an extra custodial partner, whose very strong coupling has a big effect on the width. In fact, allowing too many additional custodial partners will rapidly drive  $\Gamma_{g^1} \gtrsim M_{g^1}$ , an indication of a break-down of perturbation theory. From Eq. (16), we can infer that there can be at most four new custodial quarks whose masses are less than  $M_{g^1}/2$ .

In models with large boundary kinetic terms,  $g^1$  primarily decays into light quarks, swamping the decay into tops, and its over-all width becomes much narrower. This fact, combined with the enhancement of  $g^1$  production, allows for the possibility that one could discover  $g^1$  in the dijet mode, against the large QCD background. To explore this possibility, in Figure 8 we plot the invariant mass distribution of QCD dijets (with rough acceptance cuts  $|\eta| < 1.0$  and  $p_T > 20$  GeV to reduce the SM background). For  $M_{g^1} = 2$  or 3 TeV, we can reconstruct a peak against the dijet background with ample statistics.



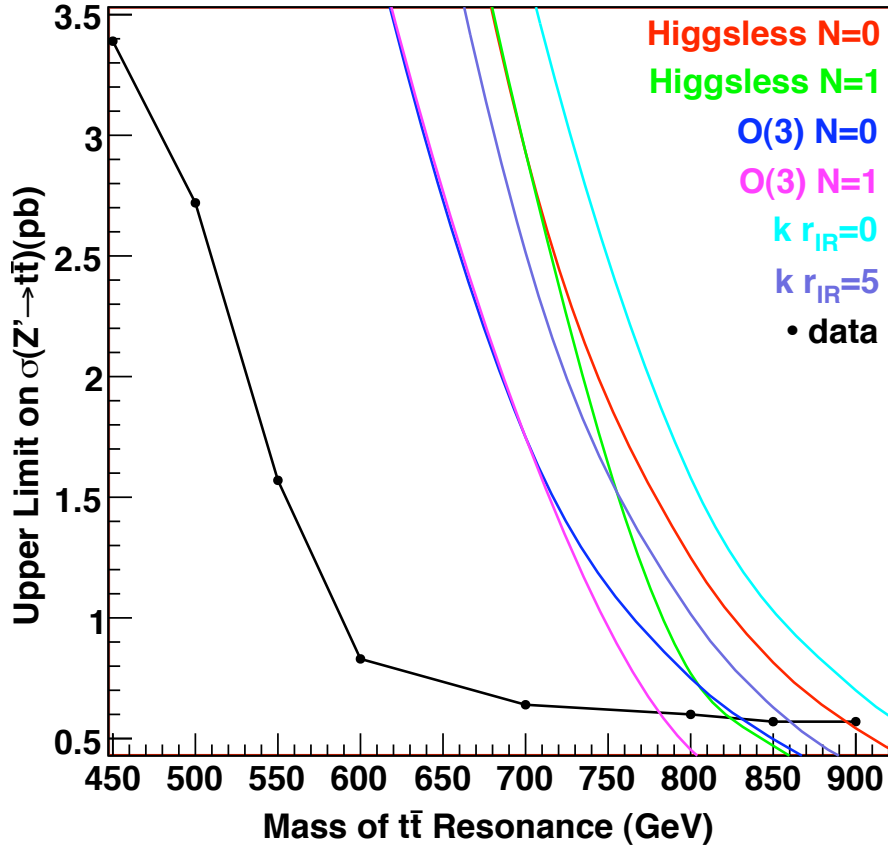


FIG. 7: Cross section for  $p\bar{p} \rightarrow g^1 \rightarrow t\bar{t}$  at the Tevatron as a function of the mass of  $g^1$ , compared with the CDF exclusion curve. The mass of custodial partners is 360 GeV.

Based on the size of the signal and background, we estimate that one could potentially discover  $g^1$  even if its mass is larger than 4 TeV in such models.

The highly chiral nature of the couplings of  $g^1$  to top, bottom, or the custodial partners may be visible as an observable [11]. The top final state is particularly promising, because the left-handed nature of the  $W$ - $t$ - $b$  interaction implies that the top decay automatically analyzes its production polarization. For example, the standard RS scenario has about 95% decays into right-polarized tops, whereas the model with  $\kappa r_{IR} = 10$  has roughly equal decays into left- and right-polarized tops, and the model with expanded custodial symmetry with  $Q_3$  localized at the IR brane has about 99% decays into left-polarized

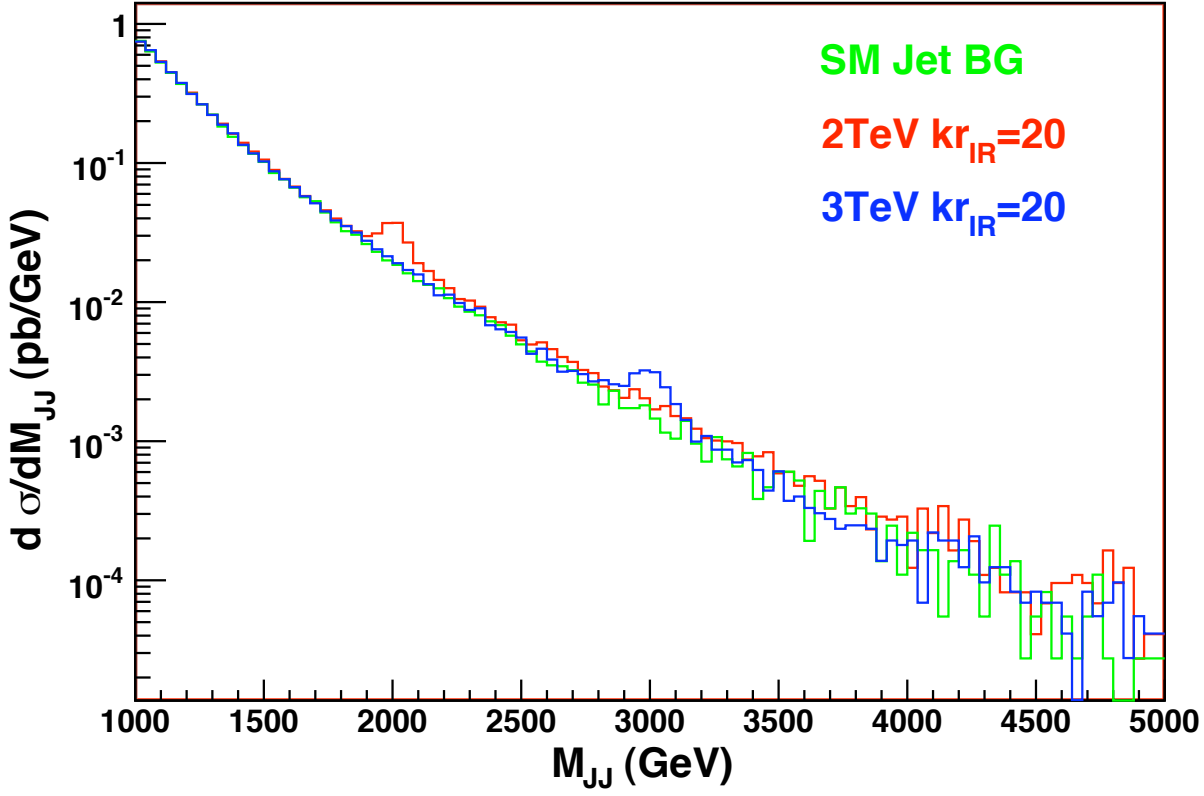


FIG. 8: Invariant mass distribution of QCD dijets coming from the KK gluon resonance in models with large brane kinetic term ( $\kappa r_{IR} = 20$ ), and the SM prediction. The cuts  $p_T > 20\text{GeV}$ ,  $|\eta| < 1.0$ , and invariant mass  $> 1\text{TeV}$  are applied.

tops.

Finally, given the large cross-sections, it is natural to ask what the current bounds from the Tevatron on anomalous top production imply for the KK gluon mass. A recent analysis from CDF [26] has set bounds on narrow resonances in the  $t\bar{t}$  invariant mass spectrum. While the analysis does not strictly apply in this case, since the KK gluon is wider than the machine resolution, the actual bound will be close to that quoted in the analysis. We have plotted this in Fig. 7, along with representative cross-sections from the models under investigation here. Note that this excludes Higgsless models with KK masses below about 850 GeV, and that includes the region favored by unitarity in  $WW$

scattering.

#### IV. INTERFERENCE

There is an intriguing feature of the fermion couplings to  $g^1$ : the sign of the coupling depends on the sign of the  $g^1$  wave function close to where the fermion is localized. As a KK mode, the  $g^1$  wave function contains a node, and changes sign from one side of the extra dimension to the other. As a result the UV fermions have a minus sign relative to the zero mode gluon coupling, while the IR fermions have a plus sign. This sign should be visible in the interference between  $s$ -channel gluon and KK-gluon production of  $t\bar{t}$ , as illustrated in Fig. 9.

To quantify this effect we propose an asymmetry parameter  $A_i$ . This parameter should be positive or negative depending on the sign of the light quark coupling and be zero in the Standard Model. We accomplish this with the definition

$$A_i = - \frac{\int dm (\frac{d\sigma}{dm} - \frac{d\sigma}{dm}_{SM}) * \Theta(m - M_{g^1})}{\int dm |\frac{d\sigma}{dm} - \frac{d\sigma}{dm}_{SM}|}. \quad (17)$$

Here  $m$  is the invariant mass in the  $t\bar{t}$  distribution and  $M_{g^1}$  is the center of the resonance. The logic of this choice is that: *i.* The SM contribution is subtracted to determine if the interference is positive or negative; *ii.* the sign of the interference changes as the resonance is crossed, hence the  $\Theta$ -function; *iii.* As is well-known, a positive sign will produce negative interference below the resonance and positive above due to the sign of the resonance propagator  $1/(s - M_{g^1}^2)$ , hence the overall minus sign. With this definition the sign of  $A_i$  will be that of the light quark coupling.

The normalization of the data with respect to the SM calculation is problematic. Since

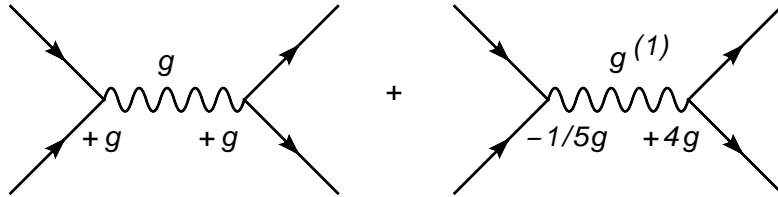


FIG. 9: Graphs that interfere allowing measurement of the sign of the light quark coupling.

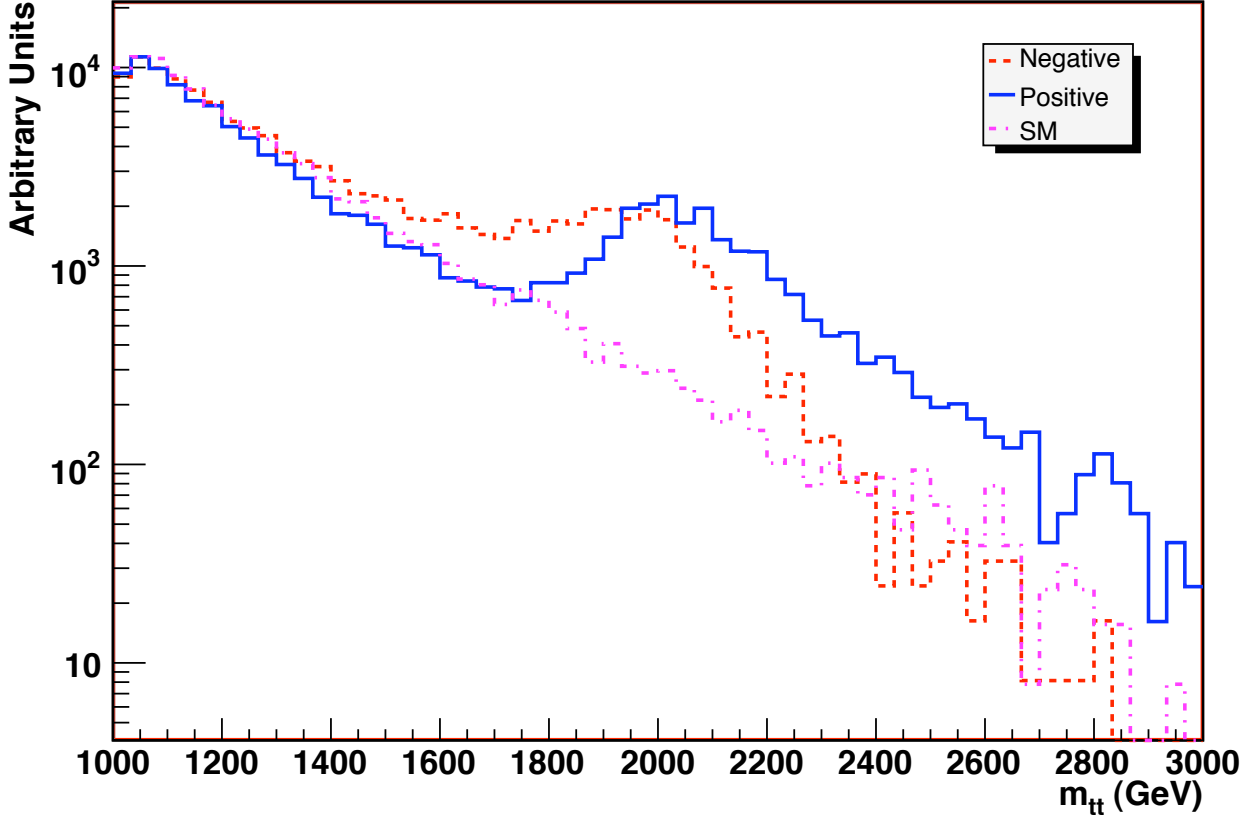


FIG. 10: Invariant mass distribution of  $pp \rightarrow t\bar{t}$  in models with positive and negative coupling to light fermions, along with the SM prediction.

the resonance will result in a much larger overall cross-section, one should not normalize to the total number of events. We choose to normalize to the lowest-mass bin used in calculating the asymmetry, which allows extraction of the normalization from data, while retaining all available information in the region near the resonance.

We present values of  $A_i$  for several masses in the basic RS model in Table II. We also show the value obtained by switching the sign of the light quark coupling. We have included a crude estimate of the smearing by shifting the value of the top and anti-top 4-momentum by a gaussian random number with width given by the ATLAS jet resolution. Since the uncertainty in top reconstruction will be dominated by the jet uncertainty this gives the correct order-of-magnitude for the smearing; we leave more refined estimates for

$g^{(1)}$ Mass	plus	minus
2 TeV	0.57	-0.44
3 TeV	0.54	-0.28
4 TeV	0.52	-0.16

TABLE II: Asymmetry parameter  $A_i$  for  $t\bar{t}$  resonances with negative (corresponding to basic RS) and positive light quark couplings.

future work. We find that the smearing makes little difference, as the resonance width is larger than the detector resolution. The results in Table II indicate that if a resonance is observed in  $t\bar{t}$  production,  $A_i$  is a promising variable to extract information about the underlying theory.

## V. CONCLUSIONS

We have investigated the structure of the KK gluon resonance in several variants of the RS model. We find that this structure contains information that will help to distinguish between models even in the absence of data from the electroweak sector. The width and branching ratios will constrain the location of the fermion zero-modes as well as the number of light KK modes into which the KK gluon can decay. In addition, the ratio of cross-sections for producing the  $g^1$  directly and in association with a  $b$ -jet will give specific information about the localization of the third generation quarks. Specifically, a large coupling to  $b\bar{b}$  will prefer a model where the  $Z \rightarrow b\bar{b}$  vertex is protected by an extended custodial symmetry. In some models, with large boundary kinetic terms, the  $g^1$  can primarily decay into dijets, and it seems promising that in such models one can discern  $g^1$  against the large QCD background up to masses somewhat larger than 4 TeV.

Finally, we find that the relative sign of the coupling to light quarks and to tops can be measured in the interference with  $s$ -channel gluon exchange. This provides an important consistency check on the overall picture of the fermion geography and the mechanism by which flavor hierarchies are realized in the fermion Yukawa couplings.

The discovery of  $g^1$  is an important first step in the discovery of RS, and further

observables such as its production rate, associated rate with bottom quarks, total width and branching ratios, and interference with SM  $t\bar{t}$  production, can yield information about the nature of the RS construction, and the parameters which describe it.

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- [1] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 4690 (1999) [arXiv:hep-th/9906064].  
L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999) [arXiv:hep-ph/9905221].
  - [2] These arguments were lucidly put together by J. March-Russell, talk at HCP 2007, Isola d'Elba, Italy.
  - [3] N. Arkani-Hamed, M. Porrati and L. Randall, JHEP **0108**, 017 (2001) [arXiv:hep-th/0012148].
  - [4] L. Randall and M. D. Schwartz, JHEP **0111**, 003 (2001) [arXiv:hep-th/0108114]; K. Agashe, A. Delgado and R. Sundrum, Annals Phys. **304**, 145 (2003) [arXiv:hep-ph/0212028]; M. Carena, A. Delgado, E. Ponton, T. M. P. Tait and C. E. M. Wagner, Phys. Rev. D **68**, 035010 (2003) [arXiv:hep-ph/0305188]; K. Agashe, R. Contino and R. Sundrum, Phys. Rev. Lett. **95**, 171804 (2005) [arXiv:hep-ph/0502222].
  - [5] S. J. Huber and Q. Shafi, Phys. Lett. B **498**, 256 (2001) [arXiv:hep-ph/0010195].
  - [6] K. Agashe and G. Servant, Phys. Rev. Lett. **93**, 231805 (2004) [arXiv:hep-ph/0403143]; K. Agashe and G. Servant, JCAP **0502**, 002 (2005) [arXiv:hep-ph/0411254].
  - [7] H. Davoudiasl, J. L. Hewett and T. G. Rizzo, Phys. Lett. B **473**, 43 (2000) [arXiv:hep-ph/9911262]; A. Pomarol, Phys. Lett. B **486**, 153 (2000) [arXiv:hep-ph/9911294]. C. Csaki, J. Erlich and J. Terning, Phys. Rev. D **66**, 064021 (2002) [arXiv:hep-ph/0203034].
  - [8] K. Agashe, A. Delgado, M. J. May and R. Sundrum, JHEP **0308**, 050 (2003) [arXiv:hep-ph/0308036].
  - [9] H. Davoudiasl, J. L. Hewett and T. G. Rizzo, Phys. Rev. D **68**, 045002 (2003) [arXiv:hep-ph/0212279]; M. Carena, E. Ponton, T. M. P. Tait and C. E. M. Wagner, Phys. Rev. D **67**, 096006 (2003) [arXiv:hep-ph/0212307]; M. Carena, A. Delgado, E. Ponton, T. M. P. Tait and C. E. M. Wagner, Phys. Rev. D **71**, 015010 (2005) [arXiv:hep-ph/0410344].
  - [10] K. Agashe, R. Contino, L. Da Rold and A. Pomarol, Phys. Lett. B **641**, 62 (2006) [arXiv:hep-ph/0605341].
  - [11] K. Agashe, A. Belyaev, T. Krupovnickas, G. Perez and J. Virzi, arXiv:hep-ph/0612015.  
B. Lillie, L. Randall and L. T. Wang, arXiv:hep-ph/0701166.

- [12] R. Contino, L. Da Rold and A. Pomarol, arXiv:hep-ph/0612048.
- [13] M. Carena, E. Ponton, J. Santiago and C. E. M. Wagner, arXiv:hep-ph/0701055.
- [14] Y. Grossman and M. Neubert, Phys. Lett. B **474**, 361 (2000) [arXiv:hep-ph/9912408].
- [15] M. Carena, T. M. P. Tait and C. E. M. Wagner, Acta Phys. Polon. B **33**, 2355 (2002) [arXiv:hep-ph/0207056].
- [16] H. Georgi, A. K. Grant and G. Hailu, Phys. Lett. B **506**, 207 (2001) [arXiv:hep-ph/0012379].
- [17] D. Choudhury, T. M. P. Tait and C. E. M. Wagner, Phys. Rev. D **65**, 053002 (2002) [arXiv:hep-ph/0109097].
- [18] K. Agashe, R. Contino and A. Pomarol, Nucl. Phys. B **719**, 165 (2005) [arXiv:hep-ph/0412089].
- [19] M. Carena, E. Ponton, J. Santiago and C. E. M. Wagner, Nucl. Phys. B **759**, 202 (2006) [arXiv:hep-ph/0607106]. A. D. Medina, N. R. Shah and C. E. M. Wagner, arXiv:0706.1281 [hep-ph].
- [20] A. Djouadi, G. Moreau and F. Richard, Nucl. Phys. B **773**, 43 (2007) [arXiv:hep-ph/0610173].
- [21] See <http://lepewwg.web.cern.ch/LEPEWWG/> for recent fits to  $S$  and  $T$ .
- [22] C. Csaki, C. Grojean, H. Murayama, L. Pilo and J. Terning, Phys. Rev. D **69**, 055006 (2004) [arXiv:hep-ph/0305237]. C. Csaki, C. Grojean, L. Pilo and J. Terning, Phys. Rev. Lett. **92**, 101802 (2004) [arXiv:hep-ph/0308038]. G. Cacciapaglia, C. Csaki, C. Grojean and J. Terning, Phys. Rev. D **70**, 075014 (2004) [arXiv:hep-ph/0401160]. G. Cacciapaglia, C. Csaki, C. Grojean and J. Terning, Phys. Rev. D **71**, 035015 (2005) [arXiv:hep-ph/0409126]. G. Cacciapaglia, C. Csaki, C. Grojean, M. Reece and J. Terning, Phys. Rev. D **72**, 095018 (2005) [arXiv:hep-ph/0505001]. C. Csaki, C. Grojean, J. Hubisz, Y. Shirman and J. Terning, Phys. Rev. D **70**, 015012 (2004) [arXiv:hep-ph/0310355]. Y. Nomura, JHEP **0311**, 050 (2003) [arXiv:hep-ph/0309189]. R. Barbieri, A. Pomarol and R. Rattazzi, Phys. Lett. B **591**, 141 (2004) [arXiv:hep-ph/0310285]. H. Davoudiasl, J. L. Hewett, B. Lillie and T. G. Rizzo, “Higgsless electroweak symmetry breaking in warped backgrounds: Constraints Phys. Rev. D **70**, 015006 (2004) [arXiv:hep-ph/0312193]. H. Davoudiasl, J. L. Hewett, B. Lillie and T. G. Rizzo, JHEP **0405**, 015 (2004) [arXiv:hep-ph/0403300]. J. L. Hewett, B. Lillie and



T. G. Rizzo, JHEP **0410**, 014 (2004) [arXiv:hep-ph/0407059]. R. S. Chivukula, E. H. Simmons, H. J. He, M. Kurachi and M. Tanabashi, “The structure of corrections to electroweak interactions in Higgsless Phys. Rev. D **70**, 075008 (2004) [arXiv:hep-ph/0406077]. R. S. Chivukula, E. H. Simmons, H. J. He, M. Kurachi and M. Tanabashi, Phys. Lett. B **603**, 210 (2004) [arXiv:hep-ph/0408262].

[23] G. Cacciapaglia, C. Csaki, G. Marandella and J. Terning, JHEP **0702**, 036 (2007) [arXiv:hep-ph/0611358]. G. Cacciapaglia, C. Csaki, G. Marandella and J. Terning, Phys. Rev. D **75**, 015003 (2007) [arXiv:hep-ph/0607146].

[24] T. Sjostrand, S. Mrenna and P. Skands, JHEP **0605**, 026 (2006) [arXiv:hep-ph/0603175].

[25] J. Alwall *et al.*, arXiv:0706.2334 [hep-ph].

[26] M. Kagan, D. Amidei, C. Cully, T. Schwarz, and M. Soderberg [http://www-cdf.fnal.gov/physics/new/top/2006/mass/mttb/pub\\_page.html](http://www-cdf.fnal.gov/physics/new/top/2006/mass/mttb/pub_page.html)

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